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COLD DENSE BARYONIC MATTER AND COMPACT STARS*

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Probing dense hadronic matter is thus far an uncharted field of physics. Here we give a brief summary of the highlights of what has been so far accomplished and what will be done in the years ahead by the World Class University III Project at Hanyang University in the endeavor to unravel and elucidate the multifacet of the cold dense baryonic matter existing in the interior of the densest visible stable object in the universe, i.e. neutron stars, strangeness stars and/or quark stars, from a modest and simplified starting point of an effective field theory modeled on the premise of QCD as well as from a gravity dual approach of hQCD. The core of the matter of our research is the possible origin of the $\sim 99\%$ of the proton mass that is to be accounted for and how the "vacuum" can be tweaked so that the source of the mass generation can be uncovered by measurements made in terrestrial as well as space laboratories. Some of the issues treated in the program concern what can be done — both theoretically and experimentally — in anticipation of what's to come for basic physics research in Korea.

Keywords: Cold dense matter; compact stars; chiral symmetry; BR scaling; dilaton limit; holographic QCD and dense matter; symmetry energy; FAIR; KoRIA.

1. The Goal

The proton which is understood to be made up of three "chiral" quarks, two up quarks (denoted as u) and one down quark (denoted as d), is the most stable hadron with the precisely measured mass $m_p = 938.272013 \pm 0.000023$ MeV. The quarks are light fermions — hence chiral — compared with the strong interaction scale ~ 100 MeV, with their mass $m_q = \frac{1}{2}(m_u + m_d) \sim 4$ MeV. How the quark masses (and also the lepton masses) arise, one of the most fundamental questions of physics,

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will be clarified once the LHC finds the Higgs particle. However the perhaps not less fundamental question is: where does the bulk of the proton mass come from?

This is a highly charged question. Take the nucleus which is the smallest object that can be understood in terms of its constituents. The mass of a nucleus of mass number A can be expressed as

$$M_A = Am_p + \delta \tag{1}$$

which is almost totally given by the sum of the proton mass with a small binding energy correction $|\delta| \lesssim 0.02m_p$. This is the case with everything that we see around us, atoms, molecules and buildings, etc. Add the mass of the constituents and that makes up the bulk. But when it comes to the proton which is a constituent of a nucleus, this picture stops. It makes no sense to write the proton mass as $m_p = 3m_q + \Delta$. This is because what corresponds to the "correction" Δ amounts to 99% of the proton mass so the question is where do most of the proton mass come from? This is one of the most fundamental questions in physics, and an attempt to answer that question leads all the way to asking what makes the densest object in the Universe, neutron star, defy the gravitational collapse to a black hole. This is the theme of the WCU3-Hanyang program.

In addressing this question, we assume our interpretation of what quantum chromodynamics (QCD) is telling us is correct: that most of the proton mass — and other hadrons' mass made up of u and d (chiral) quarks — are generated dynamically through spontaneous breaking of chiral symmetry, in other words, a sort of vacuum rearrangement. The objective of the WCU3-Hanyang program is how to "see" the making of this mass by tweaking the "vacuum structure."

The questions we pose to address this problem are the following:

(1) What happens to the hadron masses under extreme conditions such as high temperature and, in particular, high density that are believed to change the vacuum structure?
(2) How is the structure of nuclei and cold baryonic matter modified under extreme conditions?
(3) How does the modified nuclear structure influence dense matter in compact stars?
(4) What makes the neutron stars stable against gravitational collapse?

From cosmological considerations, it is not difficult to understand that a strongly interacting matter above certain temperature must be in the form of quarks and gluons, and not in hadrons. This means that going up, the temperature must "melt" the hadronic system. Indeed this understanding has been confirmed in relativistic heavy-ion collisions, where the relevant temperature has been determined to be ~ 175 MeV. Future experiments at LHC/CERN will map out the phase structure of such hot (but not so dense) matter revealing what takes place at high temperature. However the situation with high density at low temperature is totally different. It is nearly completely unknown what happens to cold hadronic system when it is squeezed to a very high density, other than that one is observing compact stars with an interior density estimated to be as high as or even higher than, say, ~ 10 times, that of nuclear matter. This is because cold dense hadronic matter is a strongly correlated system inaccessible to theoretical tools relying on standard perturbation theory, and furthermore the only reliable nonperturbative approach to QCD, i.e. lattice QCD, cannot handle, because of the famous sign problem, the density regime involved. In this paper, we describe the approach initiated by the WCU3-Hanyang team, what (little) has been so far accomplished and what remains to be done in the future. The earlier part of this development was summarized in Ref. 1.

2. Brown–Rho Scaling

The underlying framework in this project is the Brown–Rho scaling predicted in 1991 by one of the authors with G. E. Brown.[2] A basic assumption that goes into this prediction was that if one ignores the small quark mass (the process known as "the chiral limit"), the entire mass of the chiral-quark hadrons is generated by the spontaneous breakdown of chiral symmetry. Then in the limit that the number of colors N_c is taken to be large[a] for which one can justify doing tree-order calculations, the hadron masses scale in the environment of dense medium with density denoted n as

$$\frac{m_N^*}{m_N} \approx \frac{m_M^*}{m_M} \equiv \Phi(n)\,, \tag{2}$$

where the subscripts N and M stand for the nucleon (proton, neutron) and light-quark mesons (ρ, ω), respectively and the asterisk for in-medium quantity at a given density. Here $\Phi(n)$ is a scaling function that depends on density n. The property of this scaling function is a subtle issue that is not precisely understood even now after two decades since its formulation in Ref. 2. The reasoning made there is based on the notion that the unbreaking of chiral symmetry is locked to the unbreaking of spontaneously broken scale symmetry in QCD. The underlying assumption that goes into (2) which was made more precise in Ref. 3 is that the scale symmetry breaking in QCD as manifested in the trace anomaly consists — in the simplest description — of two components, spontaneous breaking and explicit breaking, with the former associated with a "soft gluon" and the latter with a "hard gluon." What figures therefore in (2) is the soft component that can be described in terms of a scalar referred to as dilaton χ which is a Goldstone boson of the spontaneously broken scale symmetry. In terms of the VEV of the dilaton in medium $\langle\chi\rangle^* \equiv \langle\chi\rangle_n$, one can associate the scaling function with the density dependence of

[a] Taking the large N_c limit is the only known nonperturbative analytical tool available so far for QCD which has been found to be reliable in certain processes for which lattice QCD calculation could be used to check, e.g. in the quenched approximation.

the "pion decay constant"[b]

$$F_\pi^* = F_\pi \frac{\langle\chi\rangle_n}{\langle\chi\rangle_0} \tag{3}$$

as

$$\Phi(n) = \frac{F_\pi^*}{F_\pi}\,. \tag{4}$$

It seems natural that this parameter appears on the r.h.s. of Eq. (2) because at low density, it is the relevant parameter connected to the order parameter of chiral symmetry, the quark condensate $\langle\bar{q}q\rangle$. That it appears linearly is a consequence of the assumption taken. It could in principle be a more involved function of it. It is important to note that the quantities in (2) are the *intrinsic or bare parameters* of the Lagrangian defined at a given density n. When $1/N_c$ corrections enter importantly via quantum loop graphs, the scaling of physical quantities such as masses will be functions of those quantities but not necessarily in the same simple form. This point is often misinterpreted in the literature causing confusion of what the evidence is for chiral symmetry restoration, partial or complete.

There are unmistakeable evidences, albeit indirect, that the scaling (2) linear in F_π^* does hold at least qualitatively or even semiquantitatively in nuclear medium up to nuclear matter density $n_0 \simeq 0.16$ fm^{-3}. There is, however, up to date no "smoking-gun" signal for its direct connection to chiral symmetry, namely, the quark condensate. It remains still a controversial issue. On the other hand, it is safe to say that contrary to what is sometimes claimed, there is no clear evidence against it either. We will return to this matter in a different context.

Naively one would think that since the (physical) pion decay constant is supposed to go to zero at high density where chiral symmetry is supposed to be restored, the mass would disappear as the pion decay constant goes to zero. However there is no theoretical argument available at present to show that the mass has to go to zero in the simple form of (2). Even if it goes to zero, it could be a complicated function of F_π^* that vanishes at the chiral restoration point. To avoid confusion, the physical pion decay constant that vanishes at chiral restoration will be denoted f_π^*. It turns out that if one implements nonlinear sigma model — the established effective field theory at low energy — with hidden local symmetry (HLS) associated with the vector mesons that are known to be relevant degrees of freedom in low-energy strong interaction dynamics[c] one finds that the r.h.s. of (2) has to be replaced as

[b]We put this quantity in the quotation mark since it is a parameter that appears in the Lagrangian which is related to the physical pion decay constant but can be identified with it only in the mean-field or tree approximation. We will make this distinction more precise later.

[c]An elegant and powerful way to incorporate the energy scale corresponding to the vector meson mass $\sim$800 MeV in the nonlinear sigma model is the hidden local symmetry approach of M. Harada and K. Yamawaki in *Phys. Rep.* **381**, 1 (2003) who showed how to consistently do chiral perturbation calculation in the presence of vector mesons and approach chiral phase transitions at high temperature and/or high density. This hidden local symmetry appears naturally in the form of infinite tower in holographic dual QCD derived from string theory, e.g. T. Sakai and S. Sugimoto, *Prog. Theor. Phys.* **113**, 843 (2005).

one approaches near the critical point — T_c or n_c — by

$$\frac{m_N^*}{m_N} \approx \frac{m_M^*}{m_M} \approx \frac{\langle \bar{q}q \rangle^*}{\langle \bar{q}q \rangle}\,. \tag{5}$$

This is a relation which should be valid *if at all* only very near the critical point with modulo an overall constant of order 1. There is at present no experimental confirmation or falsification of this relation. There are several experimental papers that claim evidences either for or against it but we consider those claims to be unfounded. A compelling argument was given in Refs. 4–6 that the experiments so far performed and analyzed by several model calculations *did not measure* the observables that are sensitive to chiral symmetry properties. The verdict is still not out, awaiting more precisely focused experiments and interpretations thereof.

It should be stressed that the hidden local symmetric approach to the scaling (5) relies on several rather strong assumptions: (1) the density driven transition is a smooth cross-over or second-order or at worst weakly first-order; (2) the matching of the one-loop HLS correlators with the operator product expansion of QCD correlators at the "chiral scale" ~ 1 GeV is reliable; and (3) the fermions, quasiquarks or baryons, in the effective Lagrangian are *the* relevant degrees of freedom. One or more of these assumptions can go wrong in several ways. There are some obvious caveats to these assumptions. Firstly, there can be a chiral-invariant mass term contributing to the nucleon or quasiquark mass, which would then invalidate both the l.h.s. and r.h.s. of (5). This possibility has been considered in terms of a parity-doublet hidden local symmetric model that allows a chiral-invariant mass term that remains nonzero in the chirally restored phase. This matter will be discussed below. Secondly in doing the matching of the correlators at a matching scale deemed optimal for both the QCD correlators and the HLS correlators, higher dimension chiral order parameters that can enter are expressed in terms of the quark condensate $\langle \bar{q}q \rangle$. But this cannot be rigorously justified when baryonic systems are simulated on crystals as discussed below. We will indeed find that at some density which will be labeled as $n_{1/2}$, $\langle \bar{q}q \rangle$ vanishes but the pion decay constant f_π^* does not, signaling that the system is still in chiral symmetry-broken phase: the nonvanishing pion decay constant must be getting contributions from higher-dimension order parameters characterizing broken chiral symmetry. Thirdly, the large N_c assumption which is the basis for the definition of the "mass" could break down, with loop corrections becoming crucial as mentioned above.

Given that cold baryonic matter at high density defies QCD-inspired approaches, it is tempting to exploit the presumed power of gravity dual models anchored on string theory to address the density regime involved. Indeed such a strategy has been found to be successful in certain processes in condensed matter and also in relativistic heavy-ion collisions that are not accessible in perturbation theory. The well-known case is the low shear-viscosity at high temperature seen in RHIC experiments which cannot be explained in perturbative QCD but can be explained in terms of the gravity-dual of $\mathcal{N} = 4$ supersymmetric theory although the latter is not QCD in the UV regime. This class of success led to the attempt to apply a holographic QCD approach to in-medium meson properties in the difficult density regime.[7] Qualitatively the scaling (2) at low density seems to be reproduced, but the high density scaling (5) was not seen at all in the model. There can be a variety of reasons for this deviation from the "gauge-theory" prediction. One mechanism unmistakenly identifiable in QCD-inspired models such as Nambu–Jona-Lasinio is the important role played by low-mass scalar mesons and $1/N_c$ corrections associated with them,[d] both of which are missing in holographic QCD. Implementing such scalars in hQCD is an open problem.

3. Half-Skyrmion Matter and New BR

The first indication that (5) may break down at some density above n_0 in dense medium was noticed when soliton configurations in nonlinear sigma model describing nucleons as skyrmions are put on an FCC crystal lattice to simulate dense matter. At present, this is the only way we know how to simulate nonperturbatively dense matter using an effective Lagrangian. Lattice gauge theory is unable to handle the density involved and all model approaches start perturbatively from the matter-free vacuum with phenomenological Lagrangians. These may be made to work up to the nuclear matter density with the help of experimental data but cannot be trusted when extrapolated to higher density.

The crystal simulation that we employ can be justified in the large N_c limit for baryonic matter at high density. What is intriguing in this approach is the observation that at some density $n_{1/2} > n_0$ (where $n_0 \approx 0.16$ fm^{-3} is the normal nuclear matter density), a skyrmion in the FCC configuration fractionizes into two half-skrymions, each carrying a half baryon charge in the BCC configuration. This feature shown in Fig. 1 is not observed in any other treatments.[e] This structure was also observed[8] with the instantons of the holographic QCD models that possess correct chiral symmetry structure put on an FCC lattice. Here what corresponds to the half-skyrmion is the dyon in the instanton structure of Ref. 8. Independently of detailed mechanisms involved, the import of this half-skyrmion-dyon identification is that the fractionized topological soliton could be a generic feature of the topological nature of the nucleon, and not exotic as has been largely thought. And it seems highly likely that it is not just an artifact of the crystal structure either. What matters is at what density the transition takes place.

What is distinctive of the half-skyrmion structure is that it arises as a sort of phase transition at a density $n_{1/2}$ lying higher than n_0, with the chiral condensate vanishing in the unit cell, $\langle\bar{q}q\rangle = 0$, but a nonvanishing pion decay constant

[d] In such models, it is seen that if N_c exceeds 3 found in nature, nuclei and nuclear matter cannot be bound. See L. Bonanno and F. Giacosa, arXiv:1102.3367 [hep-ph]. What this implices in the bulk sector is not clear and deserves to be clarified.
[e] There is a resemblance of this phase to what is called "vector symmetry" discussed by H. Georgi in *Phys. Rev. Lett.* **63**, 1917 (1989).

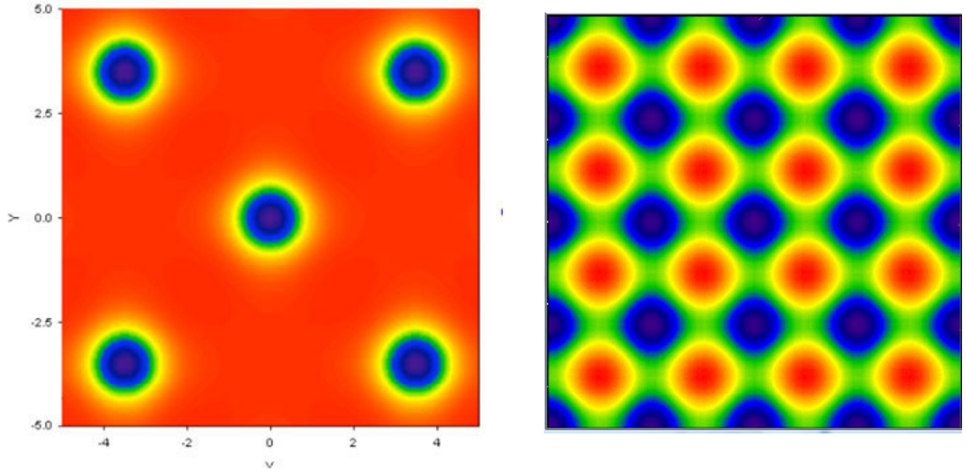

Fig. 1. (Color online) Local baryon number densities for low-density skyrmion phase (left panel) in FCC and high-density half-skyrmion phase in BCC (right panel). The red color represents the empty space. This is taken from Lee *et al.*, *Nucl. Phys. A* **723**, 427 (2003).

f_π^*. Although the quark condensate — which is the standard order parameter for chiral symmetry — is zero, the nonvanishing pion decay constant says that chiral symmetry is not actually restored. This means that it cannot be a standard order parameter. This may be symptomatic of the crystal lattice structure, which may not support chiral symmetry restoration. However it is a phase transition in that there is change in the relevant degrees of freedom from skyrmions to half-skrymions, involving also a topology change. This resembles rather "deconfined" quantum critical phenomenon in condensed matter physics which goes outside of the Landau–Ginzburg–Wilson paradigm of phase transition.[f]

There are several striking consequences in nuclear dynamics of this skyrmion-half-skyrmion phase transition, in particular in cold dense baryonic matter. When a scalar degree of freedom is incorporated via a dilaton associated with the spontaneously broken scale symmetry in QCD (in addition to the explicit breaking due to quantum anomaly, i.e. the QCD trace anomaly) into nonlinear sigma model as needed for low-energy nuclear dynamics,[3] the scaling (2) is found to hold fairly well up to $n \approx n_{1/2}$, but above $n_{1/2}$, the scaling (5) gets significantly modified. It takes the form[9]

$$\frac{m_M^*}{m_M} \approx \kappa\left(\frac{g^*}{g}\right), \tag{6}$$

$$\frac{m_N^*}{m_N} \approx \kappa\,. \tag{7}$$

[f] See T. Senthil *et al.*, *Nature* **302**, 1490 (2004).

Here g is the hidden gauge coupling constant and κ is a constant weakly dependent on density, given in the skyrmion model as F_π^*/F_π.[g] The hidden gauge coupling constant goes like $g \sim \langle\bar{q}q\rangle$ near the chiral restoration point, so in the vicinity of chiral restoration it is the hidden gauge coupling that takes over the role of chiral order parameter. Away from the critical point, how (6) scales is not known. As a whole, this modification in the scaling makes a drastic change in the structure of nuclear forces. Most notable is the change in the nuclear tensor forces.

4. Tensor Forces

The new scaling predicted by the topology change (6) and (7) can have a dramatic effect in tensor forces in dense matter for $n > n_{1/2}$. It is easy to see what happens assuming that tensor forces are mediated by the exchange of the pion and the ρ meson. Suppose that one can treat the nucleon as nonrelativistic. Then the tensor forces take the form

$$V_M^T(r) = S_M \frac{f_{\rm NM}^2}{4\pi} m_M \tau_1 \cdot \tau_2 S_{12} \times \left(\left[\frac{1}{(m_M r)^3} + \frac{1}{(m_M r)^2} + \frac{1}{3 m_M r}\right] e^{-m_M r}\right), \tag{8}$$

where $M = \pi,\ \rho,\ S_{\rho(\pi)} = +1(-1)$. The total tensor force is then the sum $V^T = V_\pi^T + V_\rho^T$ and since they come with an opposite sign, they tend to cancel.

To exhibit the consequence of the scaling modified by the half-skyrmion phase, it is convenient to express the old and new scalings as follows:

$$\left(\frac{m_N^*}{m_N}\right)_{\rm old} \approx \left(\frac{m_M^*}{m_M}\right)_{\rm old} \approx \left(\frac{f_\pi^*}{f_\pi}\right) = \Phi(n)\,, \quad \left(\frac{g^*}{g}\right) \approx 1 \quad \text{for} \quad 0 \lesssim n \lesssim n_c \tag{9}$$

and

$$\begin{aligned} &\left(\frac{m_N^*}{m_N}\right)_{\rm new} \approx \left(\frac{m_M^*}{m_M}\right)_{\rm new} \approx \left(\frac{f_\pi^*}{f_\pi}\right) = \Phi(n)\,,\\ &\left(\frac{g^*}{g}\right) \approx 1 \quad \text{for} \quad 0 \lesssim n \lesssim n_{1/2}\,,\\ &\left(\frac{m_M^*}{m_M}\right)_{\rm new} \approx \kappa\left(\frac{g^*}{g}\right) \approx \Phi'(n)\,,\\ &\left(\frac{m_N^*}{m_N}\right)_{\rm new} \approx \kappa \quad \text{for} \quad n_{1/2} \lesssim n \lesssim n_c\,. \end{aligned} \tag{10}$$

[g]The physical pion decay constant in medium or more precisely the temporal pion decay constant given by $f_\pi^{t*} = F_\pi^* + \delta F_\pi$ where δF_π stands for loop corrections, goes to zero — in the chiral limit — at the chiral restoration, but F_π^* does not.

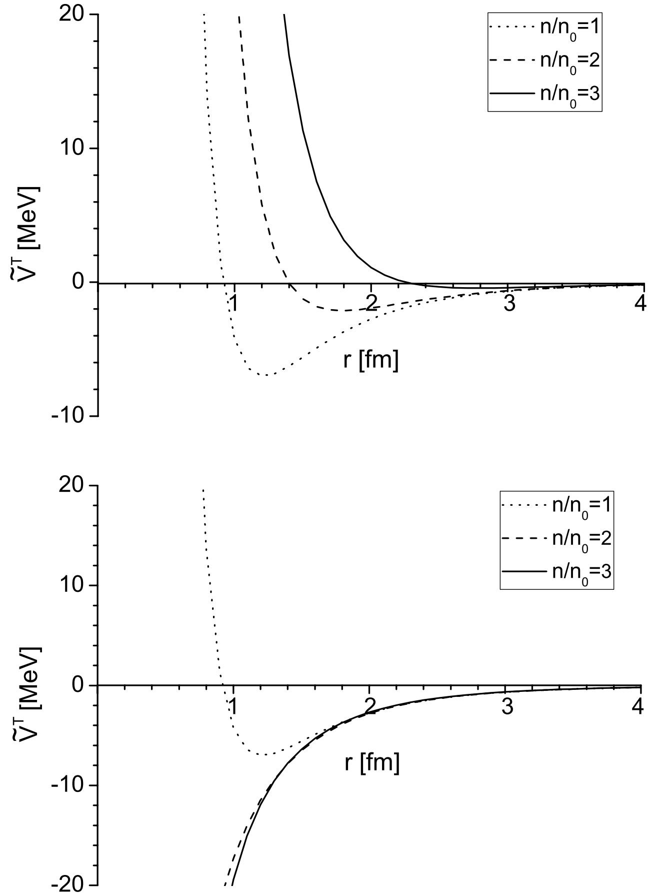

Fig. 2. Sum of π and ρ tensor forces in units of MeV as function of densities $n/n_0 = 1, 2$ and 3 with the "old scaling" (9) (upper panel) and "new scaling" (10) (lower panel).

Up to $n \approx n_0$, the two scalings are the same and constrained by experiments. Unless $n_{1/2}$ is much greater than n_0, it should be safe to take the meson and baryon scalings to be the same up to the density $n_{1/2}$. Beyond $n_{1/2}$, however, the scalings Φ and Φ' could differ. In fact, Φ' is likely to fall faster than Φ. In Fig. 2 is shown how the tensor forces would look like for a reasonable set of parameters.[9] While details depend on the parameters that we are unable to pin down precisely, the qualitative feature will remain unchanged. The new scaling predicts a tensor force structure drastically different from the old scaling above $n_{1/2}$. Specifically at $n \sim 3n_0$, the tensor force that results from the cancelation between the two is nearly zero with the old scaling, whereas with the new scaling, the ρ tensor gets strongly suppressed, leaving only the pion tensor operative. This effectively increases the tensor force for $n \gtrsim n_{1/2}$.

5. Effect of New BR on Symmetry Energy

The new scaling will affect dramatically the symmetry energy factor S in asymmetric nuclei, and of course neutron-star matter, defined by

$$E(n,x) = E(n,0) + S(n)x^2 + \cdots, \tag{11}$$

where $x = (N - P)/A$ with $A = N + P$ with N and P standing for the neutron number and the proton number, respectively. The factor S has been measured up to $n \sim n_0$ in neutron-rich nuclei but there is little information on it for $n \gtrsim n_0$ which figures importantly for the EoS of compact stars. There is no reliable theoretical tool to determine S at high density accounting for the widely diverging theoretical predictions in the literature.

It has been argued that at high densities, S is dominated by the tensor forces. A rough but good enough approximation is

$$S \propto \frac{|V^T|^2}{\bar{E}}, \tag{12}$$

where $\bar{E}$ is the average excitation energy associated with the tensor force, $\bar{E} \sim$ 200 MeV. There is kinetic energy contribution to S which is not in (12) but it is known to be small, so we shall ignore it here. With the tensor force given by the old scaling, Fig. 2, one would find that the tensor force contribution to S nearly vanish at $n \gtrsim 2n_0$. One might think that this feature agrees with what is called "supersoft" symmetry energy that vanishes at $n \sim 3n_0$.[h] However, the situation is totally different with the new scaling. There the ρ tensor gets strongly quenched for $n \gtrsim n_{1/2}$, leaving the pion tensor untouched, which means that S will increase as the net tensor force gets stronger as density increases beyond $n_{1/2}$, making the symmetry energy *stiffer again*. We will see this feature supported by the dilaton-limit fixed point discovered recently as discussed in Sec. 8. This phenomenon is expected to be highly relevant to confronting the recently measured 1.97 solar-mass neutron star. The Radioactive Ion Beam (RIB) machines — such as, e.g. the KoRIA in project in Korea — are expected to provide a valuable information on the symmetry energy S just above n_0 where experimental data are lacking.

6. The "Ice-9" Phenomenon

Another potentially important effect of the half-skyrmion phase is on deeply bound dense kaonic nuclei. It manifests itself in the attraction felt by an antikaon embedded in dense matter. The phase change can induce an enhanced attraction,[10] as shown in Fig. 3, hitherto unseen in standard nuclear theory, at $n \sim n_{1/2}$ such that a nucleus with one or more bound antikaons could become ultra-compact with the average density exceeding that of ordinary nuclear matter. Such a compact nucleus, e.g.

[h]An evidence for such a supersoft S is discussed by Z. Xiao *et al.*, *Phys. Rev. Lett.* **102**, 062502 (2009).

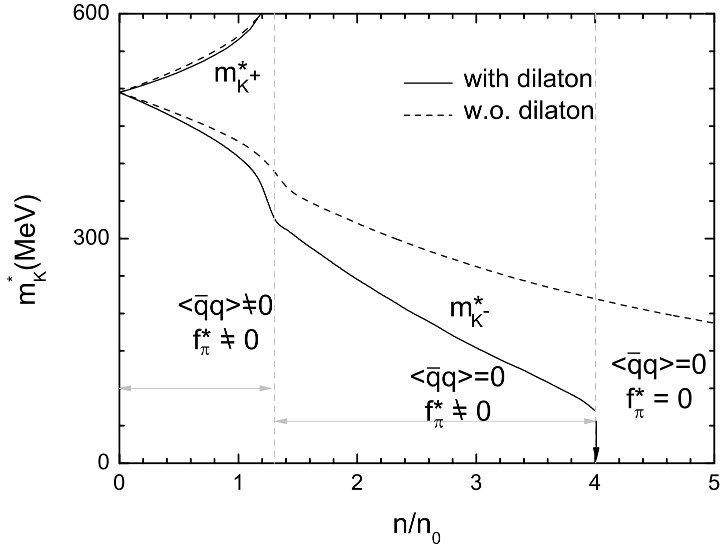

Fig. 3. $m^*_{K^\pm}$ vs n/n_0 (where $n_0 \simeq 0.16\ \mathrm{fm}^{-3}$ is the normal nuclear matter density) in dense skyrmion matter which consists of three phases: (a) $\langle\bar{q}q\rangle \neq 0$ and $f^*_\pi \neq 0$, (b) $\langle\bar{q}q\rangle = 0$ and $f^*_\pi \neq 0$ and (c) $\langle\bar{q}q\rangle = 0$ and $f^*_\pi = 0$. The role of the dilaton associated with the spontaneous breaking of scale invariance described in Sec. 2 is indicated.

K^-K^-pp, could be produced in future experiments at GSI and/or J-Parc.[i] This phenomenon has also an implication on kaon condensation in compact-star matter discussed below. The precipitous drop in the kaon mass at $n = n_{1/2}$ could trigger the formation of deeply bound antikaon-nuclear state. In conventional treatments, i.e. under normal conditions, such a process cannot occur so it may be dubbed as "ice-9" phenomenon.[j]

7. Three-Layer Compact Stars

The effect of the modification in scaling due to the topology change described in Sec. 4 is expected to influence not only the structure of finite nuclei and nuclear matter but also that of compact stars. How this modified BR scaling will affect the EoS of neutron stars is being studied in a nuclear effective field theory approach anchored on $V_{\mathrm{low}\,k}$ which is derived via Wilsonian renormalization group equations from precision nucleon–nucleon interactions.[11] This work will help provide badly needed constraints on the poorly known symmetry energy at high density.

The presence of a half-skyrmion phase at a density not too high above n_0 affects not only the tensor forces but also the antikaon spectrum. The symmetry energy factor S will also be affected by the possible role of the antikaons condensed into

[i] For a recent discussion, see T. Yamazaki *et al.*, arXiv:1106.3321 [nucl-th].

[j] Gerry Brown suggested to us the analogy between the superdense kaonic nuclei and Kurt Vonnegut's "ice-9" in *Cat's Cradle*.

the system. It has been suggested that the recently discovered high-mass neutron star PSRJ1614-2230 (with mass of $(1.97 \pm 0.04) M_\odot$) could rule out the possibility of hyperon and/or kaon condensation in compact star matter.[k] The simple reason for this assertion is that such strangeness degrees of freedom would render the EoS too soft to resist the collapse to a black hole. In the work done recently,[12] a scenario of matter composed of three phases, namely, nuclear matter at the outer layer, kaon-condensed matter in the intermediate layer and quark matter in the interior is developed into a three-layer model that could give an M vs R relation compatible with the observation.

To inject strangeness in dense baryonic matter requires treating both the hyperon degrees of freedom and the K-meson degrees of freedom simultaneously. This is a highly nonlinear and intricate matter and it is not known how to do this in a consistent way. One possibility is to start with a three-flavor hidden local symmetric Lagrangian, generate octet baryons as solitons and then treat the baryons and octet mesons on the same footing in dense medium. We are working on this problem but we are at the initial stage and have no results to show. The approach followed in Ref. 12 is to focus on the symmetry energy and consider the effect of strangeness in the symmetry energy in terms of the kaon degrees of freedom with the hyperon fields (put in by hand) in the SU(3) chiral Lagrangian integrated out.[l] The effect of the hyperons so integrated out would then be lodged in the density-dependent parameters of the kaon-nuclear Lagrangian that we will treat in the simple tree order with the multikaon interactions suppressed. There are indications from lattice simulations[m] that such multikaon interactions are not important although the effect of baryons in the multikaon interactions is neglected in the lattice calculations.

The key ingredient in our treatment for kaon condensation is the decrease of the effective mass denoted[n] as m_K^* of the negatively charged kaon K^- as density increases. The m_K^* is basically a function of m_K, ρ_n, ρ_p due to the kaon–nucleon interactions:[o]

$$m_K^* = \omega(m_K, \rho_n, \rho_p, \ldots)\,. \tag{13}$$

The density at which a neutron can decay into a proton and K^- via the weak process, $n \to p + K^-$,

$$\mu_n - \mu_p = m_K^*\,, \tag{14}$$

determines the condensed kaon amplitude ρ_t. Above the kaon condensation where m_K^* can be identified as the kaon chemical potential μ_K, the chemical equilibrium

[k] See P. B. Demorest *et al.*, *Nature* **467**, 1081 (2010).
[l] This procedure will be invalidated if there is an infrared enhancement in kaon-hyperon interactions.
[m] See W. Detmold *et al.*, *Phys. Rev. D* **78**, 054514 (2008).
[n] In this paper, m_K^* is the kaon energy in medium since we are dealing with the s-wave kaon.
[o] In this section, we will denote baryonic number density by ρ reserving n for neutron.

is reached as

$$\mu_n - \mu_p = \mu_e = \mu_\mu = \mu_K \equiv \mu\,, \tag{15}$$

where

$$\mu_n - \mu_p = 4\left(1 - 2\frac{\rho_p}{\rho}\right)S(\rho) + \Theta(K)F(K,\mu)\,, \tag{16}$$

where K stands for the kaon amplitude of kaon condensed state, i.e. $\langle K\rangle$, and $F(K,\mu)$ is a nontrivial function that depends on the neutron–proton chemical potential difference which in turn depends on kaon–nucleon interactions. It is a highly model-dependent quantity but one can see how the kaon condensation threshold depends nontrivially on the nuclear symmetry energy, represented by the symmetry energy factor $S(\rho)$.

The Hamiltonian for the nucleons and negatively charged s-wave kaons[p] involved is taken in the simple form[12]

$$\mathcal{H} = \mathcal{H}_{KN} + \mathcal{H}_{NN}\,, \tag{17}$$

where

$$\mathcal{H}_{KN} = \partial_0 K^- \partial^0 K^+ + \left[m_K^2 - \frac{n}{f^2}\Sigma_{KN}\right]K^+K^-\,, \tag{18}$$

$$\mathcal{H}_{NN} = \frac{3}{5}E_F^0\left(\frac{\rho}{\rho_0}\right)^{2/3}\rho + V(\rho) + \rho\left(1 - 2\frac{\rho_p}{\rho}\right)^2 S(\rho)\,. \tag{19}$$

Here $V(\rho)$ is the potential energy of nuclear matter as a function of density ρ and $S(\rho)$ is the symmetry energy factor as a function of nuclear density. Σ_{KN} is the KN sigma that represents the effect of the strange quark mass which is not zero (while the up and down quarks are taken as massless). It is defined as $\Sigma_{KN} \approx \frac{1}{2}(\bar{m} + m_s)\langle N|\bar{u}s + \bar{s}s|N\rangle$ and in medium, the nucleon matrix element of the bilinear quark fields will undergo medium modification. Therefore what it can be in (19) is known neither theoretically nor experimentally. Even in matter-free space, the sigma term has evolved from ~ 400 MeV to ~ 200 MeV, the latest result coming from lattice QCD measurements. In Ref. 12, this uncertainty is taken into consideration.

The potential $V(\rho)$ is fit to nuclear matter density, so considered more or less known but beyond that density, it is highly model-dependent. What we will do is to pick one convenient parametrization called MDI ("momentum-dependent interaction").[q] We shall see that while consistent within the error band with experimental constraints, they can give different results at higher density in the EoS when the strangeness and quark degrees of freedom are involved.

[p] P-wave kaons couple to the nucleon to generate hyperons that we are ignoring here.
[q] For example, the parametrization of B. A. Li *et al.*, *Phys. Rep.* **464**, 113 (2008).

In terms of the amplitude of s-wave kaon condensation, K, and the kaon chemical potential, μ_K, defined by the ansatz $K^{\pm} = Ke^{\pm i\mu t}$, $F(K,\mu)$ in Eq. (16) becomes

$$F(K,\mu) = \frac{\mu}{2f^2}K^2 + \cdots . \tag{20}$$

Following the suggestion from the lattice QCD calculation mentioned above, $\mathcal{O}(K^{2n})$ terms for $n > 1$ figuring in the ellipsis are dropped. As density increases, the kaon chemical potential $\mu(=\mu_e)$ approaches 0 at the critical density. Most of the electrons present there are converted to kaons which balance the charge neutrality of the system with the large proton fraction $x = \frac{\rho_p}{\rho} = 1/2$. Now if we consider this critical density as a phase boundary toward quark matter, then the chemical equilibrium (via confinement-deconfinement) reads

$$\mu_n - \mu_p = \mu_d - \mu_u \,, \tag{21}$$

$$\mu_{K^-} = \mu_s - \mu_u \tag{22}$$

at the phase boundary. Note that the strange quark is required at the boundary, which implies that there should be a strange quark matter for $\rho > \rho_c$. Since $\mu(=\mu_K) = 0$, we have from Eqs. (21) and (22)

$$\mu_u = \mu_d \quad \text{and} \quad \mu_s = \mu_u \,, \tag{23}$$

which gives

$$\mu_u = \mu_d = \mu_s \,. \tag{24}$$

This is the chemical potential relation for the SQM in the maselss limit. In this simple picture, the kaon condensed nuclear matter leads naturally to a strange quark matter.

For a choice of $\Sigma_{KN} \sim 260$ MeV which is compatible with current lattice results, we can have a triple-layered stellar structure consisting of NM, KNM and SQM from the outer layer to the core part as shown schematically in Fig. 4. In this work, we take a rather simple approach, namely, allow the discontinuity of density and chemical potential by assuming that NM with the density ρ_t changes into KNM with the density ρ_t' at the phase boundary defined by

$$P(\rho_t) = P(\rho_t') \,. \tag{25}$$

Assuming the strange quark matter to be in SU(3) symmetric phase in the massless limit, we have the EoS of the form

$$\epsilon_{\rm SQM} = a_4 \frac{9}{4\pi^2}\,\mu_q^4 + B = 4.83 a_4 \rho^{4/3} + B \,, \tag{26}$$

$$P_{\rm SQM} = a_4 \frac{3}{4\pi^2}\,\mu_q^4 - B = 1.61 a_4 \rho^{4/3} - B \,, \tag{27}$$

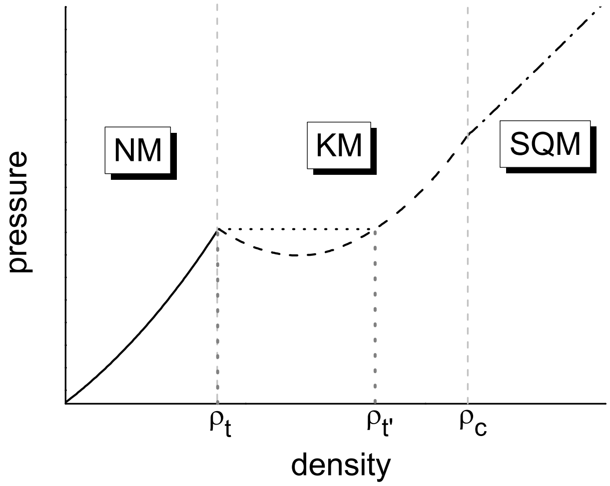

Fig. 4. The schematic phase diagram for NM-KNM-SQM.

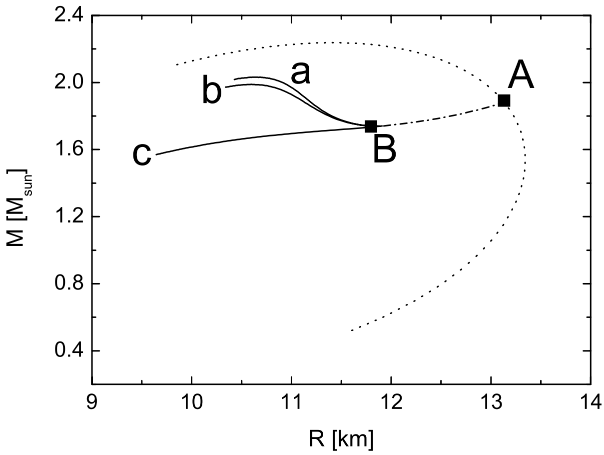

Fig. 5. The M–R sequences for LCK with $\eta = -1$. The dotted line denotes NM. The dashed-dotted line between A and B denotes the double-layered (NM-KNM) system. The solid lines, a, b and c, denote the triple-layered (NM-KNM-SQM) system with the QCD corrections with $a_4 = 0.59$, 0.62 and 1 respectively.

where B is the bag constant. Here a_4 denotes the perturbative QCD correction[r] which takes the value $a_4 \leq 1$. The equality holds for SQM without QCD corrections.

The resulting mass–radius relations that follow from the TOV equation are plotted in Fig. 5 using the parameter set, $\eta = -1$, $\Sigma_{KN} = 259$ MeV and $B^{1/4} =$

[r]See A. Chodos *et al.*, *Phys. Rev. D* **9**, 3471 (1974); J. M. Lattimer and M. Prakash, [arXiv:1012.3208 [astro-ph.SR]]; S. Weissenborn, I. Sagert, G. Pagliara, M. Hempel and J. Schaffner-Bielich, [arXiv:1102.2869 [astro-ph.HE]].

97.5 MeV. As expected, kaon condensation with higher Σ_{KN}'s would lead to smaller masses. However when the central density becomes higher, an SQM driven by kaons appears at the core. No sharp change due to the emergence of SQM is observed, which implies that the kaon driven SQM transition is a rather smooth transition in this scenario. The maximum mass of $\sim 1.99 M_\odot$ can be obtained by the gravitational instability condition at the central density $\rho = 12.3\rho_0$. With $a_4 = 0.59$, a slightly larger maximum mass, $2.03 M_\odot$, can be obtained, which is consistent with the recent observation.

What this result shows is that within the uncertainty in the parameters chosen in the model, there seems to be no great difficulty in accommodating a variety of "exotica" for the observed maximum mass of neutron stars. Much needs to be done to sharpen the arguments that figure in the theory and this is an on-going effort in the WCU-Hanyang program as listed in Sec. 11.

Given the difficulty in calculating the symmetry energy for high density in QCD-motivated approaches, it is tempting to try holographic techniques. Such a trial has been made in Ref. 13 which will be reviewed in Sec. 8. In nuclear many-body approaches, short-range correlation functions tend to dampen the short-distance repulsion in nuclear forces, in particular in the tensor forces, thereby making the symmetry energy turn over and go down at $\rho \gtrsim 4\rho_0$. The suppression of the hard-core repulsion found in the dilaton limit discussed below corroborates this feature. We will show in Sec. 9 that holographic descriptions can make a simple and unequivocal prediction that is different from what is given by nuclear many-body approaches. Whether or not this holographic QCD prediction is correct will perhaps be tested in future experiments.

8. Dilaton Limit and the Fate of Repulsive Core

The "soft" dilaton scalar figured in the mass scaling (2) in the HLS nonlinear sigma model as discussed in Sec. 2 can play an essential role in providing the attraction needed in nuclear matter that saturates at the given density. The HLS Lagrangian with baryons incorporated — which is gauge equivalent to baryonic nonlinear sigma model — cannot however describe, in mean field, what happens near the chiral phase transition. Near the transition point, the appropriate Lagrangian is presumably more like the Gell-Mann–Lévy (GML) linear sigma model. It is not known how density drives the dilatonic nonlinear sigma model to the GML-type model. But this can be achieved *by fiat* by taking what is called "dialton limit."[s]

In Ref. 14, a parity-doublet hidden local symmetry Lagrangian with the soft dilaton suitably incorporated was constructed. The objective there was to introduce the chiral-invariant fermion mass m_0 in the Lagrangian giving a nucleon mass of the form $m_N = m_0 + C$ with C that vanishes when $\langle \bar{q}q \rangle \to 0$, so making the scaling

[s]The dilaton limit used by us was discussed by S. R. Beane and U. van Kolck, *Phys. Lett. B* **328**, 137 (1994).

(5) applicable. It has been shown that one can transform the dilaton baryonic HLS Lagrangian so constructed to the GML-type Lagrangian by taking the dilaton limit

$$g_{VN} \equiv g(1 - g_V) \to 0\,, \tag{28}$$

$$g_A \to g_V\,, \tag{29}$$

$$m_N \to m_0\,, \tag{30}$$

where g is the hidden gauge coupling, $g_{V,A}$ are the vector and axial-vector couplings in the Lagrangian and g_{VN} is the effective vector-meson–nucleon coupling. The limit (30) shows that at the dilaton-limit density $n_{\rm dilaton}$, the dropping of the nucleon mass will stop at m_0. That would account for the scaling (5).

- ***Suppression of hard-core repulsion***

The most dramatic prediction in the work of Ref. 14 is that as the density $n_{\rm dilaton}$ is approached, the vector coupling disappears. A potentially important consequence is that since the repulsive core — in two-body as well as multibody forces — is generated by ω exchanges in this model, the repulsion must disappear at high density as the vector mesons decouple. This would greatly modify the EoS of compact star matter, since the density in the interior of compacts stars can reach tens of nuclear matter density. What would be needed is to determine the density at which the suppression becomes effective and how it would influence the EoS.

- ***Demise of mean-field theory at high density***

In mean-field theory, the symmetry energy goes as

$$S \sim \left(\frac{g_{\rho N}^2}{m_\rho^2}\right) n\,. \tag{31}$$

Since in the dilaton limit, the ρNN coupling goes to zero faster than the ρ mass subject to the vector manifestation (see (28)), the symmetry energy should drop rapidly as density approaches the dilaton-limit density. This is at odds with the expression (12) with the new BR. This implies that the mean-field approximation in HLS theory could go wrong at high density at least for the symmetry energy.

- ***Dilaton limit is a fixed point***

A natural question to raise is what is this "dilaton limit"? Remarkably, we find by an RGE analysis[15] that the dilaton limit exactly corresponds to the IR fixed point of HLS theory, which suggests that the system flows to that point as density is dialed. This fixed point is believed to sit slightly below the vector manifestation fixed point in HLS theory at which the (hidden) gauge coupling itself goes to zero — and the vector meson mass goes to zero.

9. Symmetry Energy from hQCD

Given that standard nuclear physics approaches give widely varying and uncontrolled results for the symmetry energy at high densities, it is highly appealing to resort to gravity/gauge duality which can in principle handle strongly-coupled phenomena like the EoS at high density. Here we give a brief summary of the effort made in that direction.[13]

Based on the treatment of dense matter in confined phase suggested in Ref. 18, a simple model for nuclear matter to strange matter transition was proposed in Ref. 19, where two D6 branes for light and intermediate mass (strange) flavors were introduced. The dense matter was introduced by uniformly distributed compact D4 branes with N_c fundamental strings attached.[t] By considering energy minimization, transition from nuclear to strange matter could be studied. To calculate the symmetry energy in nuclear matter, we consider the case where the two flavors have the same quark masses, $m_1 = m_2$. We find that the symmetry energy increases monotonically with the total charge Q which is roughly proportional to the square-root of the density.

The model we use to study the symmetry energy is the D4/D6/D6 model[19] with baryon vertices consisting of D4 branes and fundamental strings. In our approach, gluon dynamics is replaced by the gravity sourced by the N_c-colored D4 branes, and two probe D6 branes are used to describe the up and down quarks. The bare quark masses are the distances between the D4 and the two D6's in the absence of string coupling. We wrap the D4 brane on S^4 which is transverse to the original D4 brane. Due to the Chern–Simons interaction with RR-field, a U(1) gauge field is induced on the D4 brane world volume.

The source of the gauge field is interpreted as the end point of fundamental strings. Substituting the equation of motion for gauge field into the Dirac–Born–Infeld action of D4 brane with the Chern–Simons term, we get "Hamiltonian" for the D4 brane. The other endpoints of fundamental strings are attached to two D6 branes with a given number ratio. The endpoints of fundamental strings provide the source of the U(1) gauge field on the D6 brane. The two D6 branes are connected to a D4 brane by fundamental strings. Therefore, the D6 branes are pulled down and the compact D4 brane is pulled up.

As discussed in Ref. 18, the length of the fundamental strings becomes zero since the tension of the fundamental strings is always larger than that of D-branes. Finally, the position of the cusp of D6 branes should be located at the same position of the cusp of the D4 brane, ξ_c. We consider Q_1 fundamental strings attached to one D6 brane and Q_2 strings attached to another D6 brane. The force at the cusp of the D6 branes must be balanced by the pulling of the baryon vertex D4. See Fig. 6 for the final configuration of the baryon vertex and the probe flavor D6 brane.

[t]Whether this uniform distribution is a realistic picture of dense matter is an open question and remains unsettled.

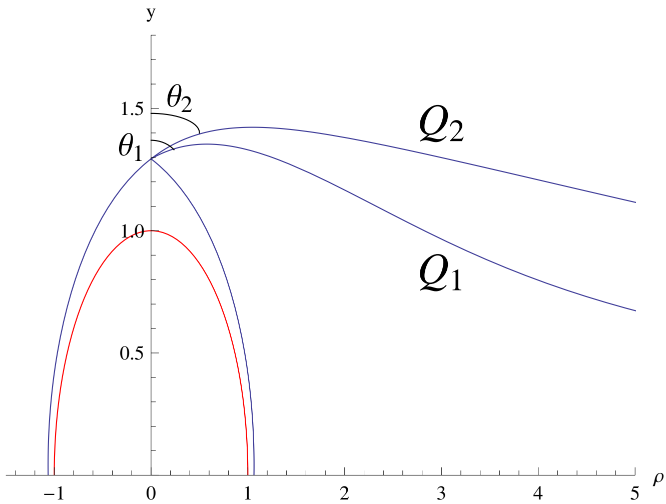

Fig. 6. (Color online) Embedding of D-branes with $\alpha \neq 0.5$. The asymptotic heights of two branes are the same ($m_1 = m_2 = 0.1$). The red curve denotes the position of U_{KK}.

For the nuclear symmetry energy, we consider only the $m_2/m_1 = 1$ case. The explicit form of the symmetry energy per nucleon can be written as

$$S_2 = \frac{2\tau_6}{N_B} \int d\rho \, \frac{\sqrt{1+\dot{y}^2}\tilde{Q}^2 \omega_+^{10/3} \rho^4}{(\tilde{Q}^2 + 4\omega_+^{8/3}\rho^4)^{3/2}} \,, \tag{32}$$

where y is the embedding solution of D6 brane with $\tilde{\alpha} = 0$. Since $N_B = Q/N_c$, the symmetry energy (32) contains an N_c factor. We need to factor this N_c out for the reason we discuss later. Our results are given in Fig. 7. Note that so far we have

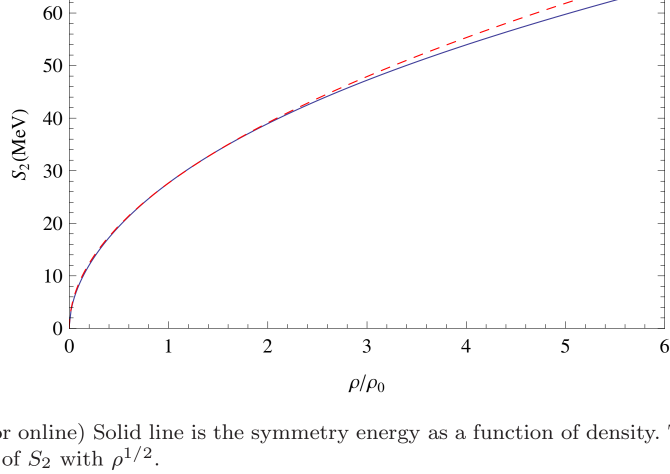

Fig. 7. (Color online) Solid line is the symmetry energy as a function of density. The dotted line is the best fit of S_2 with $\rho^{1/2}$.

used ρ for both the coordinate and the density. Hereafter ρ will denote only the density.[u] To fix the energy scale, we used the value of the 't Hooft coupling λ and the compactification scale M_{KK} determined in Ref. 20. We stress that there are two notable aspects in our results that are rather insensitive to the choice of λ and M_{KK}. One is the stiffness of the symmetry energy S_2 in the high density regime, and the other is its low density power-law behavior $S_2 \sim \rho^{1/2}$.

The power-law behavior of S_2 in the low density regime can be understood by calculating analytically in a special limit, $m_q \to \infty$ and $\rho \to 0$. In this case, the solution of the D6 brane embedding becomes trivial, $\dot{y} = 0$, and we can integrate (32) analytically to have

$$S_2 = \left(\Gamma\left(\frac{5}{4}\right)\right)^2 \sqrt{\frac{\lambda\rho_0}{2M_{KK}}}\sqrt{\frac{\rho}{\rho_0}}\,. \tag{33}$$

The current experimental result of the symmetry energy can be summarized by a fitting formula

$$S_2(\rho) = c\left(\frac{\rho}{\rho_0}\right)^\gamma \tag{34}$$

with $c \simeq 31.6$ MeV and $\gamma = 0.5$–0.7 in the low density regime, $0.3\rho_0 \le \rho \le \rho_0$.[v] With our choice of λ, M_{KK}, we obtain $\gamma \simeq 0.5$ and $c \simeq 27.7$ MeV. We note that the value of γ in our results is rather insensitive to the value of λ and M_{KK}, while the value of c depends on them.

In standard nuclear physics approaches, at very low densities $\rho \ll \rho_0$, the dominant contribution to the symmetry energy comes from the kinetic energy term which encodes the Pauli principle. This is because the kinetic contribution to the symmetry energy is $\sim \rho^{2/3}$, while the one from interactions starts from $\sim \rho^1$ due to the linear density approximation which works well at very low density. The origin of the exponent $\gamma = 2/3$ is the dispersion relation $E \sim p^2$ together with the sharp Fermi surface. In our case, the fact that $\gamma = 1/2$ suggests that either the dispersion relation is anomalous like $E \sim p^{3/2}$ or Fermi surface is fuzzy.[w] This is a striking novel prediction of hQCD, not anticipated in nuclear field theory or many-body approaches and poses an interesting future study.

10. Role of the Infinite Tower of Vector Mesons in Nucleon Structure

Although in all of the phenomena discussed above, hidden local symmetry is seen to play a significant or even crucial role, it is *not* a fundamental symmetry of QCD. It is in some sense an emergent ("hidden") symmetry arising from collective

[u] We continue to denote density by ρ as in Sec. 7.
[v] See, e.g. D. V. Shetty and S. J. Yennello, *Pramana* **75**, 259 (2010).
[w] As in condensed matter discussed by, e.g. S. S. Lee, *Phys. Rev. D* **79**, 086006 (2009) and T. Faulkner *et al.*, arXiv:0907.2694 [hep-th].

excitations of the relevant degrees of freedom, i.e. pions, quasiquarks or nucleons, etc. Up to date, only the lowest-lying vector mesons ρ, ω have been considered in the gauge symmetric framework. This is because the scale considered is the chiral scale $2\pi f_\pi \sim 1$ GeV.

Considered as an emergent symmetry, it is natural to expect that as the energy scale is increased, more hidden gauge particles could appear. In fact, the recent development of holographic QCD models indicates that an infinite tower of vector mesons can provide a more powerful framework than just limiting to the lowest excitations. An interesting question that is raised in this program is how the infinite tower of vector mesons affect the structure of dense matter.

As the first step toward this goal, we consider the nucleon electromagnetic (EM) form factors. It is well-known that the famous Sakurai vector dominance picture of 1960's works quantitatively well for the pion form factor but fails completely for the nucleon form factors. It turns out that this defect for the latter can be remedied in holographic QCD in terms of a 5D Yang–Mills action in the bulk. The vector dominance is found to work equally well for both the pion and the nucleon when the infinite tower, natural in holographic QCD, is included.

A simple but highly quantitative way to see this at low-momentum transfers is to integrate out all higher members of the tower leaving only the lowest vector mesons $V = (\rho, \omega)$ active and write the resulting form factors in terms of V in a hidden local symmetric form. The physics of the high tower will then be captured in the parameters that figure in the resulting HLS Lagrangian. Based on the consideration of chiral perturbation theory with hidden local fields, one can write, say, for the proton electric Sachs form factor for low momentum transfers $Q^2 < 1$ GeV^2 as[16]

$$G_E^p(Q^2) = \left(1 - \frac{a_E}{2}\right) + z_E \frac{Q^2}{m_\rho^2} + \frac{a_E}{2} \frac{m_\rho^2}{m_\rho^2 + Q^2} + \cdots, \tag{35}$$

where a_E and z_E are parameters given by the theory. The expression (35) can be understood as the chiral perturbation expression calculated to $\mathcal{O}(p^4)$ with the vector meson treated on the same footing as the pion. The third term with the vector-meson propagator is a specific feature characterizing the fact that the ρ meson and the pion are treated on the same footing, that is, they are of the same scale. Thus perturbative unitarity is applied both to the pion and the ρ, given in terms of the propagator which corresponds to the infinite sum of the chiral expansion. This is the unique feature of hidden local symmetry theory not shared by other treatments where the ρ mass is considered as "heavy" compared with the pion mass.

The second term is an $\mathcal{O}(p^4)$ contribution which is leading order in N_c counting. The pion loop terms contribute to the same chiral order but they are down by $1/N_c$, so can be ignored in the large N_c counting.

In the form given, the Sakurai VD corresponds to setting $a_E = 2$ and $z_E = 0$. It is given by the green curve in Fig. 8. It shows in a spectacular way what we have known since a long time, that is, it fails very badly. Holographic QCD models with infinite tower would predict what is given in blue which is not so bad for the

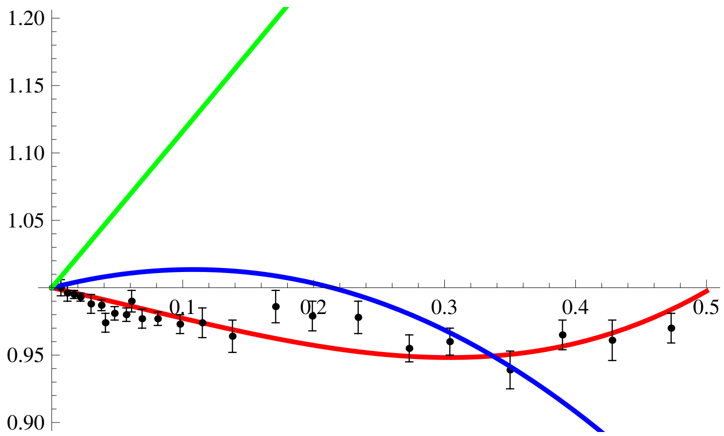

Fig. 8. (Color online) G_E^p/G_D vs Q^2. Here G_D is the well-known dipole form factor. Green: Sakurai VD; Blue: hQCD prediction; Red: best fit.

electric form factor. A best fit, which is what would be expected for the holographic QCD models if $1/N_c$ corrections were suitably taken into account, is given in red. It agrees well with Nature, with $\chi^2/\text{dof} \sim 3/2$. It shows the possibly important role of the infinite tower in the nucleon structure and more significantly, in dense baryonic matter.

It turns out[16] that the Sakai–Sugimoto holographic model which is known to be the only holographic model with the chiral symmetry of QCD in the chiral limit does not fare well with the proton magnetic moment. This is not difficult to understand. In QCD, $1/N_c$ corrections are expected to be relatively insignificant in the electric form factor but cannot be ignored in the magnetic form factor. The SS model lacks such $1/N_c$ corrections.

11. Comments and Projects

We have formulated a rather unconventional approach to dense compact-star matter starting with hidden local symmetric Lagrangian in which the dilaton degree of freedom associated with the trace anomaly of QCD and the topological solitons as baryonic degrees of freedom are incorporated. Some novel phenomena, hitherto neither observed experimentally nor predicted theoretically, have been discovered, but their validity remains up to date mostly untested. In anticipation of the forthcoming accelerators such as RIB machines (e.g. KoRIA), FAIR/GSI, J-PARC, etc. and to confront the data that will come from them, we need to sharpen the arguments that so far have been somewhat short in rigor and make the calculations more quantitative and precise. In doing so, we would like to address the following questions:

- How to systematically and accurately formulate nuclear matter, strange matter and quark matter so as to map out the phase diagram of dense baryonic matter in the strategy formulated in the program as described above? This aims to solidify the prediction — and fit to the data on the 1.97 solar mass star — made in Ref. 12.

- How does the new BR scaling affect nuclear structure, nuclear matter, neutron-rich nuclei and the EoS of neutron stars?[11] This issue is closely related to what will be measured by RIB machines (e.g. KoRIA) and FAIR/GSI.
- What is the role of the infinite tower in the short-distance properties of nuclear interactions along the line that was discussed in Ref. 16, e.g. the nucleon core size, etc. and in dense baryonic matter?[17]
- How does the infinite tower structure of skyrmions in holographic QCD affect the deeply bound kaonic nuclei and kaon condensation in dense baryonic matter?[17]
- How to exploit the power of gravity dual theory in the highly nonperturbative density regime relevant to compact stars that cannot be accessed by QCD?
- What are the specific predictions made in our approach in the density regime that can be accessed by the forthcoming RIB machines?

Closely tied to the main theme underlying the above issues, there is the possibility to probe dense matter EoS via gravity wave. This matter will be explored in the years to come.[21]

Acknowledgments

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